Three level atom optics in dipole traps and waveguides

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Abstract

An analogy is explored between a setup of three atomic traps coupled via tunneling and an internal atomic three-level system interacting with two laser fields. Within this scenario we describe a STIRAP like process which allows to move an atom between the ground states of two trapping potentials and analyze its robustness. This analogy is extended to other robust and coherent transport schemes and to systems of more than a single atom. Finally it is applied to manipulate external degrees of freedom of atomic wave packets propagating in waveguides.

1 Introduction

Exploring the wave nature of massive particles has become possible through the enormous experimental advances in the cooling of neutral atoms, ions, and molecules to temperatures where the de Broglie wavelength becomes comparable to or larger than optical wavelengths. These achievements have stimulated great interest into the field of quantum *atom* optics as an analogue of quantum optics with light [1,2]. A major objective within this field is to develop

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elements for the manipulation of the spatial wavefunction of atoms or atomic ensembles, as beam splitters, mirrors, lenses, etc. Applications are broad, ranging from a fundamental interest in probing the wave nature of particles to the manipulation of neutral atoms for implementing quantum gates and to the construction of atom interferometers for precision measurements of physical constants or as inertial sensors. In all these cases, a crucial requirement is – as in quantum optics – to preserve the coherence of the matter wave.

Of special interest to atom interferometers as well as to quantum information processing are concepts to trap or to guide atomic matter waves. For trapped atoms, the interaction with external lasers can be precisely controlled, a spreading of the wave packet can be inhibited in some or all spatial dimensions, and the effect of gravity can be compensated. Trapping and guiding of neutral atoms is usually based either on the interaction of the atom's permanent magnetic dipole moment with magnetic fields [3] or on the coupling of laser fields to the atom's induced dipole moment [4]. Arrangements of current-carrying wires [5,3], superpositions of standing light-waves with different frequencies (superlattices) [6], or appropriately shaped microlenses [7,8] allow to design and control a variety of potential shapes. Examples are Y-shaped guiding geometries to split a wave packet [9], cold atoms storage rings from guides forming a closed loop [3,5], or traps whose separation can be controlled in time [8,10,11]. As two traps are brought to a close distance and tunneling takes place, an atom initially located in one of them oscillates in a Rabi-type fashion between the two potentials. This is in close resemblance to a two-level atom interacting with a laser field, but in contrast the 'Rabi-frequency' is controlled via tuning the tunneling interaction. Such a process, if implemented correctly, is coherent as it does not introduce uncontrollable phases, and it indeed allows for a simple realization of quantum bits and quantum gates [12]. This technique however is not robust under variations of the system parameters and thus requires precise temporal control of the potentials. The same problem is of course present in optical two-level system, where for this reason a variety of robust techniques have been developed, which are based on controlling couplings in multi-level systems, and which are nowadays standard techniques in experiments.

Here we will provide a theoretical analysis of atom optical analogues to three-level techniques, especially discussing processes reminiscent of stimulated Raman adiabatic passage (STIRAP, [13]) to coherently move atoms between traps and and coherent population trapping (CPT, [14]) to create spatial superpositions of atomic wavefunctions. We will furthermore provide simulations showing that this technique is not only applicable to trapped atoms, but also to wave packets propagating in guiding potentials.

2 Time-dependent trapping potentials

To obtain an analogy between external degrees of freedom of an atom in a system of trapping potentials coupled via tunneling and an electronic three-level system coupled via the electric dipole-interaction with two laser fields, consider a linear arrangement of three atom traps. We assume strong confinement in the orthogonal directions, such that the dynamics can be restricted to the one-dimensional (1D) Hamiltonian

$$\hat{\mathcal{H}}_{\text{free}} = \int dx\, \hat{\psi}^{\dagger}(x) \left[\frac{p_x^2}{2m} + V(x,t) \right] \hat{\psi}(x) \equiv \int dx\, \hat{\psi}^{\dagger}(x) H_{\text{free}} \hat{\psi}(x), \qquad (1)$$

where $V(x,t)$ describes the potential consisting of three traps with tunable distance. At each time t, H_{free} can be diagonalized to obtain the instantaneous eigenstates. For large distance between the traps, these are localized at the centers of the traps. In the general case, states $\phi_L(x,t)$, $\phi_M(x,t)$, and $\phi_R(x,t)$, localized around the centers of the left, middle, and right trap, respectively, can be constructed from suitable combinations of the eigenstates with lowest energy. Restricting to these states, we can expand $\hat{\psi}(x) = \sum_{\alpha=L,M,R} \hat{b}_\alpha \phi(x)$ to obtain [15]

$$\mathcal{H}_{\text{free}} = -J_{LM}(t)\hat{b}_L^{\dagger} b_M - J_{MR}(t)\hat{b}_M^{\dagger} b_R - \frac{1}{2} \sum_{\alpha=L,M,R} \mu_\alpha(t)\hat{b}_\alpha^{\dagger} \hat{b}_\alpha + \text{c.c.} \qquad (2)$$

Here $J_{\alpha\beta}(t) = -\int dx\phi_\alpha^*(x,t)H_{\text{free}}(x,t)\phi_\beta(x,t)$ describes nearest-neighbor tunneling and $\mu_\alpha(t) = -\int dx\phi_\alpha^*(x,t)H_{\text{free}}(x,t)\phi_\alpha(x,t)$ are the on-site energies. Interactions of next-nearest neighbors have been neglected. Considering just a single atom and shifting the ground state energy, we arrive at the following Hamiltonian isomorphic to the Hamiltonian of a three-level system coupled by two laser fields in the rotating wave approximation [16,13]:

$$\begin{aligned} H = &-J_{LM}(t)(|L\rangle\langle M| + |M\rangle\langle L|) - J_{MR}(t)(|M\rangle\langle R| + |R\rangle\langle M|) \\ &-(\mu_M(t) - \mu_L(t))|M\rangle\langle M| - (\mu_R(t) - \mu_L(t))|R\rangle\langle R|. \end{aligned} \qquad (3)$$

The couplings $-J_{LM}$ and $-J_{MR}$ are the analogies of the Rabi frequencies of the pump and the Stokes laser, and $\mu_M - \mu_L \neq 0$ and $\mu_R - \mu_L \neq 0$ correspond to the detuning from the single- and two-photon transition, respectively (compare Fig. 1).

Optical three-level systems have been extensively analyzed. Exploiting the different possible configurations of detunings and variations of the Rabi frequencies gives rise to a large number of coherent manipulation schemes of the underlying three-level system, among them stimulated Raman adiabatic pas-
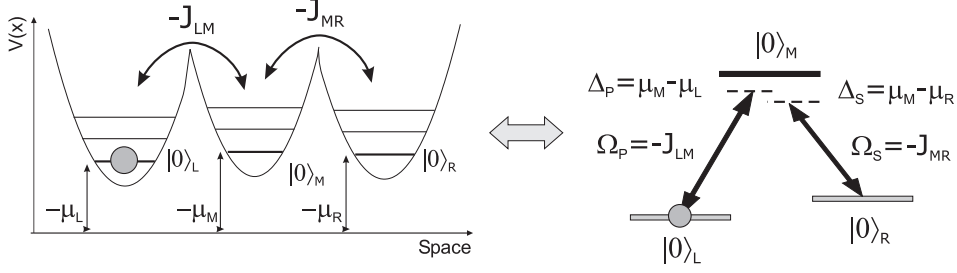
Fig. 1. Illustration of the analogy between a system of three coupled trapping potentials arranged linearly and an atomic three-level system in Λ-configuration. The tunneling matrix elements J correspond to the optical Rabi frequencies, the detunings are given by the difference between the on-site energies.

sage (STIRAP, [13]), coherent population trapping (CPT, [14]), and electromagnetically induced transparency (EIT, [17,18]). The analogy to the system of three coupled traps, as demonstrated by the Hamiltonian Eq. (3), clearly suggests to explore the application of these effects to coherently manipulate the *external* degrees of freedom of a trapped neutral atom, given the ability to control the corresponding trap parameters. Such a control should be possible in various trapping configurations as optical [7,8,19] and magnetic [3,5,11] microtraps as well as in optical lattices by exploiting superlattice techniques [6].

In the following we will especially refer to neutral atoms trapped in arrays of optical microtraps created by illuminating a set of microlenses with a red detuned laser beam [8], such that in each of the foci of the individual lenses neutral atoms can be stored by the dipole force. By illuminating the microlenses by independent laser beams under different angles, it is possible to generate various sets of traps which can be approached or separated by changing the angle between the two lasers [8]. Changing the angle allows to control the couplings between the traps, the on-site energy can be tuned by changing the laser intensities. The optical potential generated by a single laser passing through a single lens has a gaussian shape, i.e.,

$$V_{\text{trap}}(x) = -V_0 \exp\left(-\frac{1}{2V_0}m\omega_x^2 x^2\right) = -V_0 \exp\left(-\frac{\hbar\omega_x}{2V_0}(\alpha x)^2\right), \qquad (4)$$

where V_0 and ω_x are depth and frequency of the potential, respectively, m is the mass of the neutral atom and $\alpha^{-1} = \sqrt{\hbar/m\omega_x}$ is the ground state size. As several potentials of this type are superimposed, controlling the coherent and adiabatic evolution of a trapped atom is complicate as the potential depth doubles if two traps lie on top of each other. To avoid this problem, either techniques to suppress non-adiabatic excitations can be employed [20,12], or the light intensity can be reduced in an appropriate way as the traps are approached, e.g., by adding a blue detuned laser to produce a compensating extra potential [21]. Here we will focus on the later option and assume $V(x) =$

$\min[V_{\text{trap}}(x + a_L(t)), V_{\text{trap}}(x), V_{\text{trap}}(x - a_R(t))]$, where $a_L(t)$ and $a_R(t)$ fix the centers of the traps. For $m^2\omega_x^4 a_{L/R}^2/(8V_0) \ll 1$ the potential then consists of three concatenated harmonic traps.

The Hamiltonian (3) neglects contributions from non-adiabatic couplings to excited vibrational states as well as direct couplings from the left to the right trap. In what follows we will take into account the full Hamiltonian (1) through a numerical integration of the 1D Schrödinger equation to simulate the dynamics of a neutral atom in the three-trap potential.

2.1 STIRAP – robust shifting of atoms between traps

For zero detunings, one of the eigenstates of Hamiltonian (3), the *dark state*, only involves the states localized in the left and the right trap:

$$|D(\Theta)\rangle = \cos\Theta|L\rangle - \sin\Theta|R\rangle. \tag{5}$$

Here Θ is the mixing angle which depends on the couplings through $\cos\Theta = J_{LM}/J_{MR}$. To implement a robust method to move an atom from the leftmost to the rightmost trap using tunneling, the counter-intuitive STIRAP sequence can be applied: first the right and middle traps are approached and separated, and, with an appropriate delay time t_{delay}, the same sequence is used for the left and middle trap [Fig. 2 (a)]. This changes the mixing angle from $\Theta = 0$ to $\Theta = \pi/2$ [Fig. 2 (b)], and if the atom initially is located in the left trap and the process is adiabatic, then the state is at all times identical to the dark state. This moves the atom directly from $|L\rangle$ to $|R\rangle$ [Fig. 2 (c)]. The spatial wavefunctions $\langle x|D\rangle$ of the dark state for various times during the approaching process are plotted in Fig. 2 (d).

The advantage of such a STIRAP-like process, as compared to a direct transport *via* Rabi-type oscillations, is its robustness with respect to the variation of certain experimental parameters. As shown in Fig. 3 (a), the scheme works for a large range of delay times t_{Delay} and minimum distances d_{min}. A similar robustness is found for variations of, *e.g.*, the duration t_r of the approaching/separation process and the time t_i for which the traps are kept at constant distance, the only requirements being the adiabaticity of the process and the order of approaching and separating the traps. In an experimental realization, certainly a shaking of the centers of the trapping potentials provides an important source of decoherence. It can, e.g., be caused by a mechanical vibration of the microlenses, or by variations in the laser phases for optical lattices. Here we anticipate a periodic variation of the distance of the traps with frequency well below the trapping frequency, namely $\omega_{\text{Shake}} = 10^{-2}\omega_x$. As
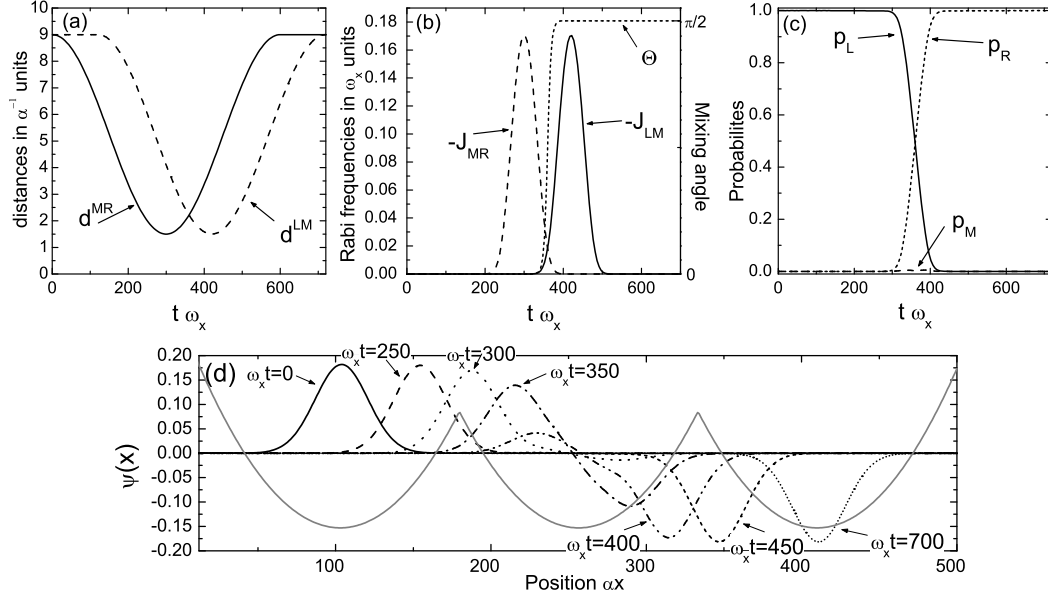
Fig. 2. (a) Approaching sequence for a STIRAP-like process, (b) the evolution of the tunneling Rabi frequencies $-J_{LM}(t)$ and $-J_{MR}(t)$ between the left and the middle and the middle and the right trap, calculated from the tunneling splitting energy of two traps, together with the mixing angle $\Theta(t) = \arctan(J_{LM}(t)/J_{MR}(t))$, and (c) the corresponding ground state populations; The parameters are $d_{max}^{LM}\alpha = d_{max}^{MR}\alpha = 9$ (maximal distances between traps), $d_{min}^{LM}\alpha = d_{min}^{MR}\alpha = 1.5$ (minimal distances), $t_r^{LM}\omega_x = t_r^{MR}\omega_x = 300$ (time used to approach/separate the traps), $t_i^{LM}\omega_x = t_i^{MR}\omega_x = 0$ (time at which the traps are at the minimal distance, and $t_{delay}\omega_x = 120$ (delay between the the approaching processes). (d) The spatial wavefunction $\langle x|D\rangle$ of the dark state for various times. The gray line gives the potential at $\omega_x t = 0$.

Fig. 3 (b) shows, the transport efficiency is not significantly degraded even for shaking amplitudes on the order of a few percent of the minimal trap distance if the delay time is appropriately chosen.

A parameter difficult to control in an experiment is the exact horizontal alignment of the traps. If a slight tilt is present in the potential, then gravity will change the relative depth of the potential minima. In this case, as has been reported in Bose-Einstein condensation in a double-trap potential [22], after a sufficiently adiabatic evolution, the atom(s) will eventually be found in the trap with lower energy. To allow for a transport to the desired state, the evolution should be explicitly non-adiabatic [23]. Here we will show that the STIRAP-like transport is within a large range not affected by gravity. To this aim we add a potential

$$\Delta V_{\text{tilt}}(x) = \gamma\hbar\omega_x\alpha x \tag{6}$$

stemming from gravity to the dipolar trapping potential $V(x)$. The parameter γ determines the slope of the ramp, and we will use $\gamma > 0$, such that the right
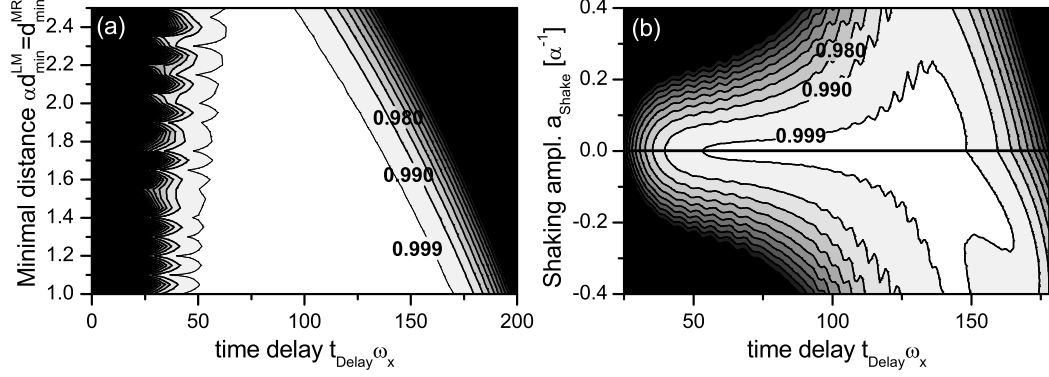
Fig. 3. Robustness of the atom optics version of STIRAP, i.e., the transfer efficiency from $|L\rangle$ and $|R\rangle$, measured by the population $\rho_R = |\langle R|\psi(t)\rangle|^2$. All parameters not varied in the figures are as in Fig. 2. In (a) the delay time t_{delay} between the two approaches (horizontal axis) and the minimal distances between traps (vertical axis) are modified. (b) shows the transfer efficiency as a function of t_{Delay} (horizontal axis) and of the amplitude of a shaking a_{Shake} in the positions of the outer traps (vertical axis) with $\omega_{Shake} = 10^{-2}\omega_x$. For $a_{\text{Shake}} > 0$ the shaking of the outer traps is in phase, for $a_{\text{Shake}} < 0$ it is out of phase by π.
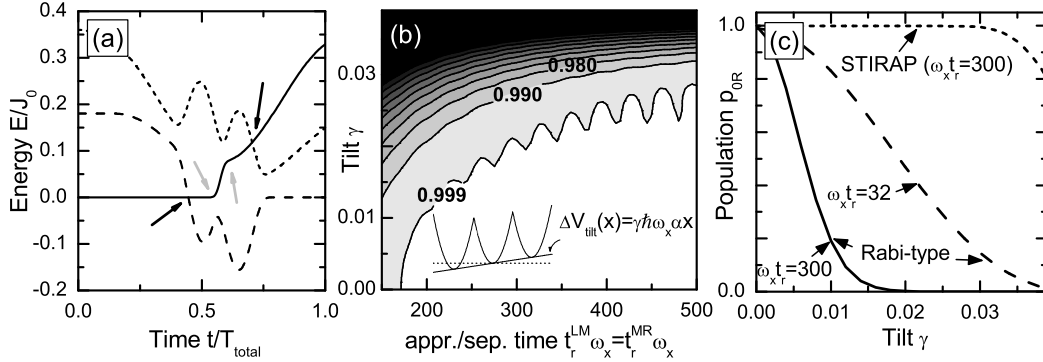

Fig. 4. Transfer efficiency for an additional tilted potential $V_{\text{tilt}}(x) = \gamma\ \hbar\omega_x\ \alpha x$, parametrized by γ. (a) Temporal variation of the energy levels obtained from a diagonalization of Eq. (3) for the parameters as in Fig. 2 and $\gamma = 0.02$. Black arrows correspond to diabatic crossings, gray arrows indicate points where the probability of non-adiabatic transitions is larger. In (b) the transfer efficiency is plotted as a function of the time t_r needed to approach and separate the traps and of the tilt; in (c) the dependence of the transfer efficiency on the tilt of the potentials is compared for STIRAP for the parameters from (b) with $\omega_x t_r = 300$ and for the transfer *via* Rabi-type oscillations between two traps for $\omega_x t_r = 300$ ($\omega_x t_i = 12$ such that full population transfer occurs for $\gamma = 0$) and $\omega_x t_r = 32$ ($\omega_x t_i = 25$).

trap is shifted up in energy with respect to the left one. For the parameters of our simulations, a value of $\gamma = 10^{-2}$ corresponds to a difference in the potential energy of $3 \cdot 10^{-2}\ \hbar\omega_x$ between the outer traps at the minimal distance. For $\gamma \ll 1$ such a tilt affects only the on-site energies $-\mu_{\alpha\beta}$ in the Hamiltonian (3). In the picture of an optical Λ-system this corresponds to a shift from the one- as well as from the two-photon resonance. In this case there exist no

adiabatic path from $|L\rangle$ to $|R\rangle$ [24]. This is exemplified in Fig. 4 (a), which shows the energies of the three eigenstates of the Hamiltonian Eq. (3) for the parameters of the STIRAP sequence from Fig. 2 and $\gamma = 2 \times 10^{-2}$. To obtain the transport from the left to the right trap, the process has to be designed as a combination of diabatic (black arrows) and adiabatic (grey arrows) processes. However, the conditions to obtain a diabatic crossing at the points indicated by the black arrows are usually fulfilled, such that the transfer efficiency is dominated by the adiabaticity requirement. For this reason, for a given γ the fidelity improves if the approaching and separating of the the traps is made slower as can be seen from Fig. 2. As should be stressed again, this is in contrast to the Rabi-type transport between two traps, where a faster process has larger fidelity of the atom ending up in the initially empty trap. Fig. 3 (d) shows a comparison of the two schemes: next to the transfer efficiency for STIRAP for $\omega_x t_r = 300$ as a function of the tile γ two curves for Rabi-type oscillations are shown for a slow ($\omega_x t_r = 300$) and for a fast ($\omega_x t_r = 32$) approaching process.

2.2 CPT-like and EIT-like effects

From the isomorphism between the Hamiltonians, it is obvious that also other processes from three-level optics can be exploited here. The approach sequence can be modified to create spatial superposition states with maximum atomic coherence by evolving the mixing angle from $\Theta = 0$ to $\Theta = \pi/4$, corresponding to a delayed approach of the left trap, but to symmetric separation. This process, reminiscent of coherent population trapping (CPT), is similarly robust to the variation of parameters as the STIRAP process, *as long as the symmetry of the separation process is maintained*. A detailed analysis can be found in [25], where also an EIT-like process is described.

2.3 Effects of atom–atom interaction

To arrive at the Hamiltonian of Eq. (3), the system has been simplified through a restriction to the lowest energy levels and to only a single atom. Our numerical simulations, which took into account excited states, show that for adiabatic processes and the potentials considered here the former simplification is justified and allows to reproduce effects known from Λ systems. Excited states however can be exploited explicitly by starting with the atom in a state different from the ground state [25]. On the other hand, also the restriction to a single particle can be released. For sufficiently low temperatures, interaction between bosonic atoms is dominated by s-wave scattering, and restricting
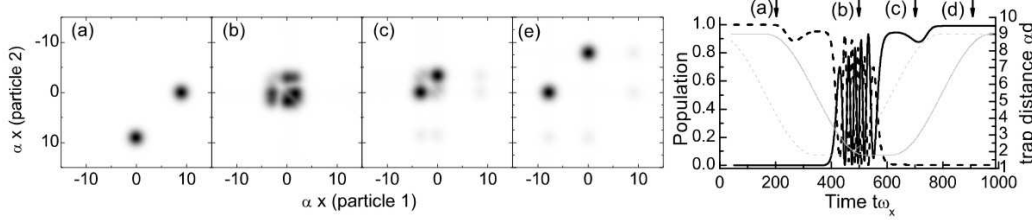
Fig. 5. Robust and coherent transport of a hole from the left trap to the right trap in a system of three traps filled with two bosonic ^{87}Rb atoms with scattering length $a_t = 106a_0$. Shown are plots of the two particle probabilities $|\psi(x_1, x_2)|^2$ at four different times (a-d) as indicated by the arrows in (e). The initial state is $|\psi(t_{\text{init}})\rangle = |0_M\rangle|0_R\rangle + |0_R\rangle|0_M\rangle$, the parameters are $t_r^{LM}\omega_x = t_r^{MR}\omega_x = 350$, $t_i^{LM}\omega_x = t_i^{MR}\omega_x = 100$, $t_{\text{delay}}\omega_x = 180$, $d_{max}^{LM}\alpha = d_{max}^{MR}\alpha = 9$, and $d_{min}^{LM}\alpha = d_{min}^{MR}\alpha = 1.5$.

again to states ϕ_α, the Hamiltonian describing the system is modified as follows [26,27]:

$$\mathcal{H} = \mathcal{H}_{Free} + \frac{1}{2} \sum_{\alpha=L,M,R} U_\alpha \hat{b}_\alpha^\dagger \hat{b}_\alpha^\dagger \hat{b}_\alpha \hat{b}_\alpha. \tag{7}$$

Here $U_\alpha = 4\pi\hbar\tilde{a}_{\text{sc}} \int dx |\phi_\alpha(x)|^4$, where $\tilde{a}_{\text{sc}}$ is the 1D scattering length which can be changed *via* changing the orthogonal confinement or exploiting a Feshbach resonance. If $|U_\alpha|$ is sufficiently large to separate sectors in energy space with different particle number within one trap, then interaction allows to move coherently more than one atom at once, or to generate 'Schrödinger cat'-like states through a CPT sequence. On the other hand, starting from a system of three traps and two bosonic atoms initially in different traps allows to coherently and robustly transport the 'hole', i.e., the empty site. An optical analogue of such a system has been studied in [28], where coherent population trapping has been analyzed for two electrons with aligned spins in a three-level system. Also in this case a dark state exists which can be interpreted as the dark state of a 'hole'. A similar effect can be achieved in the atom optical system. As an example, Fig. 5 demonstrates the corresponding STIRAP process which moves the hole between the outer traps.

3 Manipulation of matter waves in guiding structures

In the previous part we have, in close analogy to the three-level processes for internal atomic states, manipulated the external wavefunction of a trapped atom by a temporal variation of the coupling between traps. We will demonstrate that similar methods allow to manipulate an atomic wave packet propagating in an appropriately designed *fixed* guiding structure. We will assume a system of three waveguides oriented in the direction of the y axes, with y-dependent distances (see Fig. 6 (a) for an example), and a corresponding Hamiltonian

$H_{\text{free}} = (p_x^2 + p_y^2)/2m + V(x, y)$. Now, instead of considering the eigenstates of the 1D potential for each fixed time t, we can compute eigenstates at each position y, and as before combine the states with lowest energy to states $\phi_\alpha(x, y)$, $\alpha \in \{L, M, R\}$ localized around the center of each waveguide. The full wavefunction can then be decomposed as $\psi(x, y, t) = \sum_\alpha c_\alpha(y, t)\phi_\alpha(x, y)$, and inserting this expression into the Schrödinger equation gives the following equation for the evolution of the coefficients $c_\alpha(y, t)$ [9]:

$$i\hbar \frac{\partial c_\alpha}{\partial t} = -\frac{\hbar^2}{2m}\frac{\partial^2 c_\alpha}{\partial y^2} + \sum_{\beta=L,M,R}(H_{\alpha\beta} + \frac{\hbar^2}{m}P_{\alpha\beta})c_\beta + \frac{\hbar^2}{2m}\sum_{\beta=L,M,R}K_{\alpha\beta}\frac{\partial c_\beta}{\partial y}. \quad (8)$$

Here $H_{\alpha\beta}(y) = \int dx\, \phi_\alpha^*(x, y)(p_x^2/2m + V(x, y))\phi_\beta(x, y)$ are the Hamiltonian matrix elements for fixed position y and $K_{\alpha\beta}(y) = -\int dx\, \phi_\alpha^*(x, y)\partial_y\phi_\beta(x, y)$ and $P_{\alpha\beta}(y) = -\int dx\, \phi_\alpha^*(x, y)\partial_y^2\phi_\beta(x, y)$ are the kinetic and potential couplings, respectively.

For time-dependent trapping potentials, the STIRAP or CPT sequence was induced by the counterintuitive temporal ordering of the approaching and separation processes of the traps. Similarly, in the case of waveguides we will apply such sequences in space. Thus, to obtain a STIRAP–like transport with the atom initially located in the left arm, first the right guide is approached to the middle one, then, with an appropriate delay, the tunneling is also switched on between the middle and the left guide. Finally, tunneling is turned off in the same order: first between the right and the middle and then between the middle and the left guide. For a CPT-like process to split an atomic wave packet coherently between two waveguides, the approaching sequence has the same counterintuitive order, but the separation is symmetric, see Fig. 6 (a). Now the additional coupling terms in Eq. (8) make the evolution more complex. Expanding c_α into plane waves with momentum $\hbar k$ leads to a diagonal k^2-proportional term which accounts for broadening of the wave packet and to a term proportional to $k\, K_{\alpha\beta}$ which induces velocity-dependent couplings between the waveguides. Furthermore, a velocity-independent modification of the couplings is introduced through the potential couplings $P_{\alpha\beta}$.

To take again into account further effects beyond this illustrative approximation in order to evaluate the performance of such processes we have numerically integrated the full 2D Schrödinger equation. We assume to initially have an atomic wave packet located in the left arm, with a gaussian profile in the direction of the waveguide corresponding to mean momentum $\langle k_y \rangle$ and momentum spread $\Delta k_y = k_r$ ($k_r = \sqrt{2m\omega_r/\hbar}$; for the simulations $\omega_r = \omega_x/6$). In the transverse direction the wave packet corresponds to the ground state of the potential. Fig. 6 (b-d) shows an example of the time evolution in a structure which generates a splitting of the wave packet through a CPT–like configura-
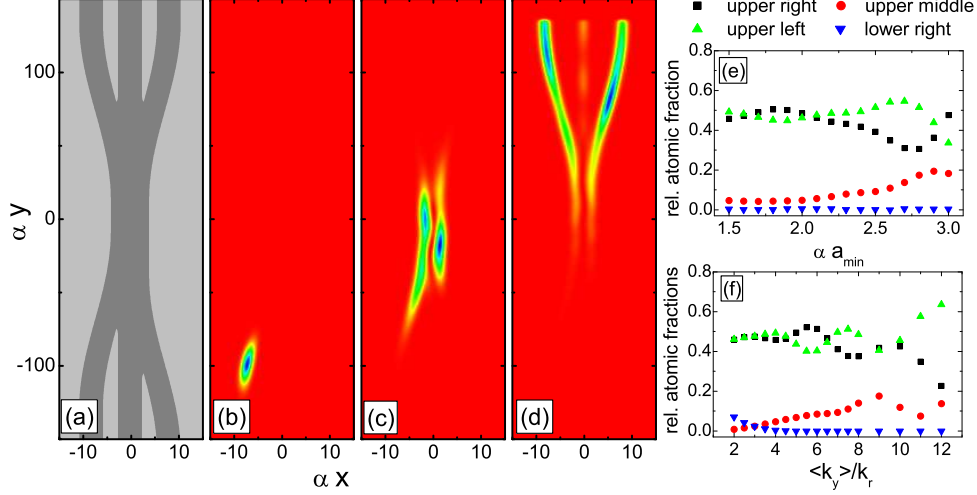
Fig. 6. (a) Contour plot of the waveguide potential used to split the wave packet incident in the left arm into a superposition of packets traveling in the left and the right arm using a CPT–like process. The contour line shown corresponds to a distance of $3/2$ ground state widths away from the centers of the waveguides. (b)–(d) show density plots of $|\psi(x,y)|^2$ at times $\omega_x t = 20$, $\omega_x t = 60$, $\omega_x t = 120$ for a wave packet with mean momentum $\langle k_y \rangle = 3.5 k_r$ and initial width $\Delta k_y = k_r$. The minimal distance between waveguides is $\alpha d_{\min} = 1.5$. (e) The relative atomic fractions leaving the setup through the upper and lower exits of the structure as the minimal distance between waveguides is modified. (f) As (e), but for different mean velocities $\langle k_y \rangle$.

tion. During the evolution, the wave packet strongly broadens in the direction of propagation, but is nevertheless nearly equally split between the left and right outgoing arms with a negligible amount of reflection. This splitting is again relatively robust with respect to the parameters describing the potential, provided the symmetry of the splitting is maintained. Fig. 6 (d) shows the change of the atomic fractions in the exit ports of the setup as the minimal distance of the waveguides is varied. The process is not as perfect as its counterpart in traps due to the additional couplings present here. Especially the velocity-dependent couplings modifying the desired CPT–like process play an important role: the larger the mean velocity, the stronger the deviation from the equal splitting, as can be seen from Fig. 6 (e).

4   Conclusions

We have studied the manipulation of the external wavefunction of an atom in a dipole potential consisting of three traps whose coupling can be changed as a function of time. As we have demonstrated, such a system, restricted to the lowest eigenstates, constitutes an analogon to the extensively studied system

of three internal atomic states coupled via two external laser fields. This allows to apply concepts as STIRAP, CPT, or EIT to coherently and robustly manipulate the external atomic wavefunction. Such processes are of potential interest, e.g., to move around atomic quantum bits or to create superposition states for interferometry. In particular, we have analyzed the robustness of a STIRAP–like process allowing transport of the atom between trapping potentials, and we have also shown that coherent processes are possible for several interacting atoms.

As a different setup, we have studied atomic wave packets propagating in waveguide potentials, where the time dependence of the trap distances is replaced by a spatial variation of the distance between waveguides. Due to additional, partially velocity-dependent couplings, the evolution is more involved and the transport or splitting processes are not as clean as in the case of traps. Still, a stronger robustness as for schemes relying only on Rabi-type tunneling between traps can be achieved, as we have exemplified through demonstrating the coherent splitting of a wave packet between two arms, a scheme interesting for, e.g, interferometry.

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